

A Gauge-Mediation Model of Dynamical SUSY Breaking with a Wide Range of the Gravitino Mass

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Abstract

We provide a gauge-mediation model of dynamical supersymmetry breaking where the gravitino mass takes a value in a wide range from 0.1 eV to 10 GeV. The lower mass region of order 100 keV or less deserves experimental and cosmological interests. The vacuum of our model is a true one without runaway instability, which may be desirable in a cosmological evolution of the universe.

1 Introduction

If it is caused by nonperturbative dynamics of some gauge interaction, low-energy supersymmetry (SUSY) breaking may provide a solution to the hierarchy problem. In recent years, many mechanisms for dynamical SUSY breaking (DSB) have been found and many realistic gauge-mediation models have been constructed where the DSB effects are transferred to the standard-model sector through the known gauge interactions [1]. The gauge-mediation models of DSB are very attractive not only because they provide calculable models for the mass spectra of superparticles in the SUSY standard model (SSM), but also because they can naturally solve the flavor-changing neutral current problem and the CP violation problem in the SSM.

Most of the gauge-mediation models considered so far [1] predict a gravitino mass $m_{3/2} \gtrsim 1$ MeV and few models are known to have $m_{3/2} \lesssim 100$ keV [2, 3]. However, models with the lighter gravitino $m_{3/2} \lesssim 100$ keV are interesting experimentally since the lightest superparticle beside the gravitino can decay into the gravitino emitting photons inside the detector, which may be testable in future experiments [4]. Also it has been pointed out [5] that the lighter gravitino is cosmologically favored if string moduli exist at the electroweak scale.

The purpose of this paper is to provide a novel gauge-mediation model where the gravitino mass takes a value in a wide range $0.1 \text{ eV} \lesssim m_{3/2} \lesssim 10 \text{ GeV}$.¹ The vacuum of our model is a true one without runaway instability [6, 7], which may be desirable in a cosmological evolution of the universe.

2 The Model

Our model is based on two sectors of simple gauge theories. One is an SU(2) gauge theory with 4 doublet chiral superfields Q_i , where i denotes a flavor index ($i = 1, \dots, 4$). The other is an SO(10) gauge theory with chiral superfields H in the 10-dimensional representation and ψ in the 16-dimensional representation. The model consists of an

¹The upperbound of the gravitino mass $m_{3/2} \lesssim 10 \text{ GeV}$ is determined so that the flavor-changing neutral currents induced by superparticle loops are sufficiently suppressed.

SU(2) and an SO(10) DSB models [8, 9] combined by means of gauge singlets through a superpotential.

The superpotential is given by

$$W = \lambda_Z Z_a (QQ)_a + \lambda_Y Y (QQ) + \frac{1}{2} f_Y S Y^2 + \frac{1}{2} f_H S H^2 + k S q \bar{q}, \quad (1)$$

where $(QQ)_a$ ($a = 1, \dots, 5$) and (QQ) denote suitable linear combinations [2, 10] of the gauge invariants $Q_i Q_j$ and Z_a , Y , S , q and \bar{q} are all singlets of $SU(2) \times SO(10)$.² It is straightforward to introduce messenger leptons in addition to the messenger quarks q and \bar{q} [1], but here we omit them for simplicity of presentation. The term $k S q \bar{q}$ transmits the DSB effects to the messenger quarks, which is a minimal setting of gauge mediation and predicts a definite mass spectrum of superparticles in the SSM [1].

In the following sections, we analyze the model by means of effective field theories. We adopt the coupling constants of order one and a phase convention that they are positive.

3 The Effective Superpotential

We first integrate out the sector of the SU(2) gauge theory. Assuming that the coupling λ_Z is relatively large, we can also integrate out the singlets Z_a along with the doublets Q_i to obtain [2, 10]

$$W_{eff} = \lambda_Y \Lambda^2 Y + \frac{1}{2} f_Y S Y^2 + \frac{1}{2} f_H S H^2 + k S q \bar{q}, \quad (2)$$

where Λ denotes a dynamical scale of the SU(2) gauge interaction.

Without the SO(10) gauge interaction, the model would possess a SUSY-invariant vacuum with $S \rightarrow \infty$. This runaway behavior is circumvented by the SO(10) gauge interaction, which gives SUSY-breaking effects when the 10-dimensional representation H is decoupled by a mass $|f_H S|$ and only the 16-dimensional representation ψ is left over

²Any symmetry which allows all the terms in Eq. (1) cannot ensure the absence of a term $\frac{1}{2} f_Z S Z_a^2$ which destabilizes the SUSY-breaking vacuum, so that we have to forbid this term by hand. However, the superpotential can be made natural if we take the couplings between singlets and Q 's to be $\lambda_Z Z_a (QQ)_a + \lambda_Y Y (QQ)^2 / M^2$, where M is the gravitational scale. Then, the gauge group acting on Q 's should be $SP(2N)$ ($N > 1$, $SP(2) \simeq SU(2)$) instead of $SU(2)$, and Λ^2 should be replaced by Λ^4 / M^2 in the following discussions [11].

[7, 9]. Conversely, if the SU(2) gauge interaction were absent, the model would have a SUSY-invariant vacuum with a mass $|f_H S|$ of the 10-dimensional representation H equal to zero [9]. In contrast to the models in Ref. [6, 7], the present model avoids an undesirable vacuum with $S = 0$ due to the sector of the SU(2) gauge theory (see the next section).

The contribution of the SO(10) sector to the effective superpotential is added [9] to yield

$$W_{eff} = \lambda_Y \Lambda^2 Y + \frac{1}{2} f_Y S Y^2 + \frac{1}{2} f_H S H^2 + k S q \bar{q} + \frac{2^{\frac{1}{5}} \tilde{\Lambda}^{\frac{21}{5}}}{(\psi \psi H)^{\frac{2}{5}}}, \quad (3)$$

where $\tilde{\Lambda}$ denotes a dynamical scale of the SO(10) gauge interaction.³ We may write this effective superpotential as

$$W_{eff} = \lambda_Y \Lambda^2 Y + \frac{1}{2} f_Y S Y^2 + f_H S H^+ H^- + k S q \bar{q} + \frac{\tilde{\Lambda}^{\frac{21}{5}}}{(\chi^2 H^+)^{\frac{2}{5}}} \quad (4)$$

in terms of variables H_+ , H_- , χ defined in Ref. [9] with a D -flatness condition

$$|H^+|^2 - |H^-|^2 - \frac{1}{2} |\chi|^2 = 0. \quad (5)$$

Here and henceforth, we denote the scalar component of a chiral superfield by the same symbol as the corresponding superfield.

4 The Effective Potential

In the following analysis, we concern ourselves with the regime $\tilde{\Lambda} \lesssim \sqrt{\lambda_Y} \Lambda$ of our interest.

4.1 The $|f_H S| \ll \tilde{\Lambda}$ region

Let the effective Kähler potential around the minimum of the potential under the fixed value of S be approximately canonical in the variables H_+ , H_- , χ and the singlets in the effective theory. This treatment seems adequate for small $|f_H S|$, which results in large $|H|$ compared to $\tilde{\Lambda}$ [9].

The effective potential around the minimum under the fixed S is given by

$$V_{eff} \simeq |\lambda_Y \Lambda^2 + f_Y S Y|^2 + \left| \frac{1}{2} f_Y Y^2 + f_H H^+ H^- + k q \bar{q} \right|^2 + |k S q|^2 + |k S \bar{q}|^2$$

³The SO(10) sector can be replaced by an SU(5) DSB model [9] without affecting our conclusions.

$$+|f_H S H^+|^2 + \left| f_H S H^- - \frac{2}{5} \frac{\tilde{\Lambda}^{\frac{21}{5}}}{(\chi^2 H^+)^{\frac{7}{5}}} \chi^2 \right|^2 + \left| \frac{4}{5} \frac{\tilde{\Lambda}^{\frac{21}{5}}}{(\chi^2 H^+)^{\frac{7}{5}}} \chi H^+ \right|^2 \quad (6)$$

$$= |\lambda_Y \Lambda^2 + Y'|^2 + |q'|^2 + |\bar{q}'|^2 + \frac{1}{|S|^4} \left| \frac{1}{2} f_Y^{-1} Y'^2 + f_H^{-1} H'^+ H'^- + k^{-1} q' \bar{q}' \right|^2 \\ + |H'^+|^2 + \left| H'^- - \frac{2}{5} \frac{(f_H S)^{\frac{11}{5}} \tilde{\Lambda}^{\frac{21}{5}}}{(\chi'^2 H'^+)^{\frac{7}{5}}} \chi'^2 \right|^2 + \left| \frac{4}{5} \frac{(f_H S)^{\frac{11}{5}} \tilde{\Lambda}^{\frac{21}{5}}}{(\chi'^2 H'^+)^{\frac{7}{5}}} \chi' H'^+ \right|^2 \quad (7)$$

with the D -flatness condition $|H'^+|^2 - |H'^-|^2 - |\chi'|^2/2 = 0$, where the primed variables are defined by

$$Y' = f_Y S Y, \quad q' = k S q, \quad \bar{q}' = k S \bar{q}, \quad H'^{\pm} = f_H S H^{\pm}, \quad \chi' = f_H S \chi. \quad (8)$$

When k is relatively large, $q = \bar{q} = 0$ is energetically favored. Then the effective potential around the minimum under the fixed S may be approximated by

$$V_{eff} \simeq |\lambda_Y \Lambda^2 + Y'|^2 + \frac{1}{|S|^4} \left| \frac{1}{2} f_Y^{-1} |Y'|^2 - f_H^{-1} |H'^-|^2 \right|^2 + 2 |H'^-|^2 \quad (9)$$

with the aid of $|H'^+|^2 = |H'^-|^2 + |\chi'|^2/2$, since $|f_H S| \ll \tilde{\Lambda} \lesssim \sqrt{\lambda_Y} \Lambda$. Hence the minimum of the potential with the fixed S is given by

$$V_{eff} \simeq \frac{f_H \lambda_Y^2 \Lambda^4}{f_H + f_Y} - f_H^2 |S|^4, \quad (10)$$

where the F term of S turns out to be

$$|F_S| \simeq f_H |S|^2. \quad (11)$$

The potential Eq. (10) shows that $|S|$ tends to run away from the origin.

4.2 The $|f_H S| \gg \tilde{\Lambda}$ region

When $|f_H S|$ is large, we may integrate out the sector of the $SO(10)$ gauge theory to obtain an effective potential of the form

$$V_{eff} \simeq |\lambda_Y \Lambda^2 + f_Y S Y|^2 + \left| \frac{1}{2} f_Y Y^2 + k q \bar{q} \right|^2 + |k S q|^2 + |k S \bar{q}|^2 + c |f_H S|^{\frac{2}{11}} \tilde{\Lambda}^{\frac{42}{11}}, \quad (12)$$

where c is a positive constant of order one [7]. When k is relatively large, $q = \bar{q} = 0$ is again energetically favored. Then the effective potential is given by

$$V_{eff} \simeq |\lambda_Y \Lambda^2 + f_Y SY|^2 + \left| \frac{1}{2} f_Y Y^2 \right|^2 + c |f_H S|^{\frac{2}{11}} \tilde{\Lambda}^{\frac{42}{11}}. \quad (13)$$

For $|\sqrt{f_Y} S| \ll \sqrt{\lambda_Y} \Lambda$, the minimum of the potential under the fixed S is given by

$$V_{eff} \simeq \lambda_Y^2 \Lambda^4 - \frac{3}{2^{\frac{2}{3}}} \left(\lambda_Y \Lambda^2 \left| \sqrt{f_Y} S \right| \right)^{\frac{4}{3}} + c |f_H S|^{\frac{2}{11}} \tilde{\Lambda}^{\frac{42}{11}}, \quad (14)$$

where the F term of S turns out to be

$$|F_S| \simeq \frac{1}{2^{\frac{1}{3}}} \left(\lambda_Y \Lambda^2 \left| \sqrt{f_Y} S \right| \right)^{\frac{2}{3}}. \quad (15)$$

For $|\sqrt{f_Y} S| \gg \sqrt{\lambda_Y} \Lambda$, the minimum of the potential under the fixed S is given by

$$V_{eff} \simeq \left| \frac{\lambda_Y^2 \Lambda^4}{2 f_Y S^2} \right|^2 + c |f_H S|^{\frac{2}{11}} \tilde{\Lambda}^{\frac{42}{11}}, \quad (16)$$

where the F term of S turns out to be

$$|F_S| \simeq \left| \frac{\lambda_Y^2 \Lambda^4}{2 f_Y S^2} \right|. \quad (17)$$

We now see that the runaway of $|S|$ observed above is stopped by the second term [7] in the potential Eq. (16).

5 The Vacuum

The analysis in the previous section shows that the effective potential V_{eff} has the true minimum at

$$|S| \simeq \left(\frac{11}{2c} \right)^{\frac{11}{46}} \left(\frac{\lambda_Y^2 \Lambda^4}{f_Y f_H^{\frac{1}{11}} \tilde{\Lambda}^{\frac{21}{11}}} \right)^{\frac{11}{23}} \gg \sqrt{f_Y^{-1} \lambda_Y} \Lambda, \quad (18)$$

when $\tilde{\Lambda} \ll \sqrt{\lambda_Y} \Lambda$. The gluino and gravitino masses [1] are given by

$$m_{1/2} \simeq \frac{\alpha_3}{4\pi} \left| \frac{F_S}{S} \right| \simeq \frac{\alpha_3}{4\pi} \left| \frac{\lambda_Y^2 \Lambda^4}{2 f_Y S^3} \right|, \quad m_{3/2} \simeq \sqrt{\frac{V_{eff}}{3M^2}}, \quad (19)$$

where $\alpha_3 = g_3^2/4\pi$ denotes the color gauge coupling, V_{eff} is given by Eq. (16) and M is the gravitational scale $\simeq 2.4 \times 10^{18}$ GeV. Since the couplings λ_Y, f_Y, f_H and c are of order one, the gluino and gravitino masses are given by $m_{1/2} \sim (\alpha_3/4\pi)(\tilde{\Lambda}/\Lambda)^{63/23}\Lambda$ and $m_{3/2} \sim (\tilde{\Lambda}/\Lambda)^{42/23}(\Lambda^2/\sqrt{3}M)$. We certainly have a wide range of the gravitino mass by means of two independent scales Λ and $\tilde{\Lambda}$. For instance, $m_{1/2} \simeq 200$ GeV and $m_{3/2} \simeq 100$ keV are realized when $\Lambda \simeq 6 \times 10^8$ GeV and $\tilde{\Lambda} \simeq 10^7$ GeV.

In the extreme case $\tilde{\Lambda} \sim \sqrt{\lambda_Y}\Lambda$, the analysis in the previous section implies that the gluino mass $m_{1/2}$ and gravitino mass $m_{3/2}$ are of order $(\alpha_3/4\pi)\tilde{\Lambda}$ and of order $\tilde{\Lambda}^2/\sqrt{3}M$, respectively. For the gluino mass $m_{1/2} \simeq 200$ GeV we get $\tilde{\Lambda} \simeq 2 \times 10^4$ GeV and a very light gravitino $m_{3/2} \simeq 0.1$ eV.

6 Conclusion

We have constructed a gauge-mediation model of dynamical supersymmetry breaking whose vacuum is a true one without runaway instability [6, 7]. Beside the SSM sector, the model consists of an SU(2) and an SO(10) DSB models [8, 9] combined by means of gauge singlets through a superpotential Eq. (1).

The gravitino mass is determined by the ratio of dynamical scales of the SU(2) and SO(10) gauge interactions and takes a value in a wide range $0.1 \text{ eV} \lesssim m_{3/2} \lesssim 10 \text{ GeV}$. In particular, a light mass of order 100 keV or less deserves experimental and cosmological interests. The present model is unique in the point that it may accommodate $0.1 \text{ eV} \lesssim m_{3/2} \lesssim 1 \text{ keV}$ with the simplest messenger term $kSq\bar{q}$ of gauge-mediated SUSY breaking. It is well known [12] that this mass region for the gravitino has no cosmological problem of gravitino overproduction.

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